

CENTER VORTICES AND COLOUR CONFINEMENT IN LATTICE QCD*

ROMAN BERTLE, MANFRIED FABER

Institut für Kernphysik, Technische Universität Wien, A-1040 Vienna, Austria
E-mail: bertle@kph.tuwien.ac.at, faber@kph.tuwien.ac.at

JEFF GREENSITE

*Physics and Astronomy Department, San Francisco State University,
San Francisco, CA 94117, USA, and
Theory Group, Lawrence Berkeley National Laboratory, Berkeley, CA 84720, USA*
E-mail: greensit@stars.sfsu.edu

ŠTEFAN OLEJNÍK

Institute of Physics, Slovak Academy of Sciences, SK-842 28 Bratislava, Slovakia
E-mail: fyziolej@savba.sk

We review lattice evidence showing that center-vortex condensation is a serious candidate for the mechanism of colour confinement in quantum chromodynamics.

1 Introduction

Any model that contends for the true theory of colour confinement in $SU(N)$ gauge theory has to be able to provide an explanation of or an answer to the following facts and questions:

- All non-zero N -ality colour charges are confined.
- All zero N -ality charges are screened at *large* distances.
- At *intermediate* distances, potentials between higher representation charges rise approximately linearly, and the corresponding string tensions are proportional to the eigenvalues of the quadratic Casimir operator of the representation. (We dubbed this phenomenon *Casimir scaling*.)
- A deconfinement transition occurs at high temperatures or high densities of hadronic matter. What is the nature of the transition? What happens to confining configurations at the transition, what symmetry gets restored or broken?
- What is the relation to other phenomena, like chiral symmetry breaking?

In this talk I will argue that confining configurations that have a chance to successfully address the above issues, are *center vortices*. I will mostly review

*Plenary talk presented by Š. Olejník. His work was supported in part by the “Action Austria–Slovakia: Cooperation in Science and Education” (Project No. 30s12) and the Slovak Grant Agency for Science (Grant VEGA No. 2/7119/2000).

recent results of our group obtained in lattice investigations of the confinement problem. Other aspects of the vortex picture were covered at this Conference in the plenary talk of Reinhardt¹ (see also short contributions of Alexandru², Engelhardt³, and Langfeld⁴).

The vortex-condensation model of confinement stems from late seventies, and was formulated first by 't Hooft⁵ and developed by many authors. Let me briefly recall the original idea. To characterize phases of the gauge theory, Wilson⁶ suggested to use the loop operator

$$A(C) = \text{Tr} \left[P \exp \left(ig \oint_C \mathbf{A}_\mu dx^\mu \right) \right]. \quad (1)$$

In a pure gauge theory, confinement is signalled by the leading behaviour of this quantity, namely

$$W(C) = \langle A(C) \rangle \sim \begin{cases} e^{-\sigma \mathcal{A}(C)} & \dots & \text{confinement,} \\ e^{-\mu \mathcal{P}(C)} & \dots & \text{deconfinement,} \end{cases} \quad (2)$$

where $\mathcal{A}(C)$ is the area of the minimal surface spanned by the loop, $\mathcal{P}(C)$ its perimeter. Wilson's loop operator measures the *magnetic* flux through C and creates *electric* flux along C .

't Hooft⁵ suggested another operator, measuring the *electric* flux through some other loop C' and creating *magnetic* flux along C' . It was defined through commutation relation with $A(C)$ (n is the *linking number* between C and C'):

$$A(C)B(C') = B(C')A(C) \exp(2\pi i n/N). \quad (3)$$

The expected behaviour of 't Hooft's operator is opposite to the Wilson loop:

$$\langle B(C') \rangle \sim \begin{cases} e^{-\mu \mathcal{P}(C')} & \dots & \text{confinement,} \\ e^{-\sigma \mathcal{A}(C')} & \dots & \text{deconfinement.} \end{cases} \quad (4)$$

The effect of the 't Hooft loop operator $B(C')$ on physical fields is a gauge transformation which is *singular* on the curve C' . It creates a thin vortex of magnetic flux (line-like in 3D, surface-like in 4D). To avoid infinite energy, the singularity along the loop C' has to be smoothened: vortices acquire a *core* of certain thickness. *Many things in this talk will be related in some way to the thickness of the vortex core.*

The essence of the vortex model of confinement is that the QCD vacuum is filled with closed magnetic vortices that have the topology of tubes (in 3D) or surfaces (in 4D). Thick vortices condense in the vacuum. A simple argument can explain the area law for the Wilson loop through fluctuations in the number of vortices linked to the loop.

It is a sort of paradox that this idea has for years not been subjected to extensive numerical tests in lattice simulations. However, in the last couple of years the idea returned to the stage and got into the focus of attention. This turn was due to the discovery of center dominance in maximal center gauge^{7,8}, and the following accumulation of evidence in favour of the picture^{9–14}.

The outline of my talk is the following: In Section 2 I will briefly sketch our recent strong-coupling calculation¹⁵, which shows that center vortices are stable saddlepoint configurations of a long-range effective action that can be derived from the usual lattice Wilson action. Then (Section 3), I will introduce our method to detect center vortices that is based on center projection in maximal center gauge, and review some interesting results obtained in numerical simulations. Section 4 then addresses a question, raised recently by Bornyakov et al.¹⁶, how physical results depend on the number of gauge copies used in iterative MCG fixing. Finally, I will argue in Section 5 that the observed Casimir scaling can be accommodated with the vortex picture, and that the thickness of the vortex core is vital for the explanation of the scaling behaviour.

2 Confining Configurations of the Long-Range Effective Action

There are various ways to argue for the center of the gauge group being important for the confinement mechanism. The asymptotic string tension for charges in the representation r depends only on its N -ality (i.e. Z_N charge); the deconfinement transition is associated with spontaneous breaking of the global Z_N symmetry; etc. We have recently provided another argument: if one calculates an effective action at larger distances (and strong coupling) starting from the usual Wilson action, its stable saddlepoints are thick center vortices.

The calculation proceeds in the following way¹⁵: We start from the pure gauge theory at strong coupling on a fine lattice (links on the fine lattice are denoted generically as U), and try to calculate the effective action on a coarse lattice with links V , where the coarse lattice spacing is L times the original U -lattice spacing. The idea is to derive an effective action where the *leading* contributions to any Wilson loop on the V -lattice are obtained from a *local* action. To achieve this, we integrate over all links on the U -lattice except \tilde{U} -links on small 2-cubes around sites of the V -lattice, as shown in Fig. 1.^a

This defines effective action \tilde{S}_L :

$$Z = \int [DV] \int \prod_{\substack{l \in \\ \text{2-cubes}}} [d\tilde{U}_l] \left\{ \int \prod_{\substack{l'' \notin \\ \text{2-cubes}}} [dU_{l'']} \prod_{l'} \delta \left[V_l^\dagger (UU..U)_{l'} - I \right] e^{S_W[U]} \right\}$$

^aWe work in 3 dimensions. The extension to 4D should be straightforward.

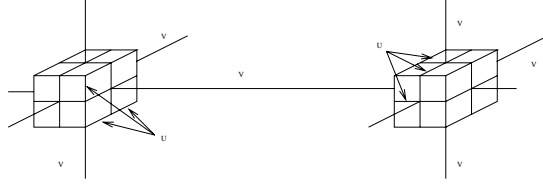


Figure 1: Unintegrated links on the coarse V -lattice.

$$\equiv \int [DV] \int \prod_{l \in 2\text{-cubes}} [d\tilde{U}_l] \exp \left(\tilde{S}_L[V, \tilde{U}] \right). \quad (5)$$

Instead of \tilde{U}_l we introduce in the 2-cubes group-valued plaquette variables h, g (Fig. 2)¹⁷. After integrating out the g variables, the result is (displaying only terms of low order in both h and β):

$$\begin{aligned} Z \approx & \int [DV] [Dh] \prod_K \left\{ 1 + 2 \left(\frac{\beta}{4} \right)^3 \sum_{c \in K} \chi_{\frac{1}{2}}[(hhh)_c] \right. \\ & \left. 2 \left(\frac{\beta}{4} \right)^4 \sum_{\text{adj. } c_1 c_2 \in K} \chi_{\frac{1}{2}}[(hhh)_{c_1} (hhh)_{c_2}] + \dots \right\} \\ & \times \exp \left[\frac{\beta}{2} \sum \text{Tr}[h] + 2 \left(\frac{\beta}{4} \right)^{4(L-2)} \sum_{l'} f^{ijkl} \text{Tr}[h_{ij}^\dagger V_{l'} h_{kl}^\dagger V_{l'}^\dagger] \right. \\ & \left. + 2 \left(\frac{\beta}{4} \right)^{L^2} \sum_{P'} \text{Tr}[V V V^\dagger V^\dagger] \right]. \quad (6) \end{aligned}$$

This resembles an adjoint-Higgs Lagrangian, with an $SU(2)$ gauge field V_μ coupled to 24 “matter” fields h in the adjoint representation. Note that

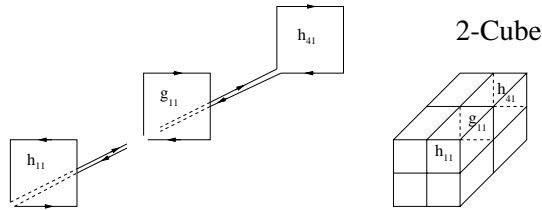


Figure 2: Plaquette variables on the 2-cubes.

for large L , the “Higgs” potential term is much larger than the “kinetic” and pure-gauge (V -plaquette) terms, so the h -fields fluctuate almost independent of V_μ .

We then make a unitary gauge fixing of the h -fields (breaking $SU(2)$ to Z_2), integrate out the remaining h d.o.f., and arrive at:

$$\begin{aligned} S_{eff}[V] &\approx S_{link}[V, \langle h \rangle_h] + S_{plaq}[V] \\ &= 2 \left(\frac{\beta}{4} \right)^{4(L-2)} \sum_{l'} f_{l'}^{ijkl} \text{Tr} \left[\langle h_{ij}^\dagger \rangle_h V_{l'} \langle h_{kl}^\dagger \rangle_h V_{l'}^\dagger \right] \\ &\quad + 2 \left(\frac{\beta}{4} \right)^{L^2} \sum_{P'} \text{Tr}[V V V^\dagger V^\dagger]. \end{aligned} \quad (7)$$

Let us now look for saddlepoints. S_{link} is maximized by

$$V_\mu(\vec{n}) = Z_\mu(\vec{n}) \times g(\vec{n}) g^\dagger(\vec{n} + \mu), \quad Z_\mu = \pm \mathbf{1}, \quad (8)$$

where $g(\vec{n}) g^\dagger(\vec{n} + \mu)$ is fixed by the particular unitary gauge choice, while S_{plaq} is maximized if $ZZZZ = +\mathbf{1}$. This is the unitary gauge *ground state*.

One can create a *thin center vortex* on this state by a discontinuous gauge transformation, e.g.

$$Z_y(\vec{n}) = \begin{cases} -\mathbf{1} & \dots & n_1 \geq 2, n_2 = 1, \\ +\mathbf{1} & \dots & \text{otherwise,} \end{cases} \quad Z_x(\vec{n}) = Z_z(\vec{n}) = \mathbf{1}. \quad (9)$$

This configuration is stationary: S_{link} is still a maximum (insensitive to center), S_{plaq} is extremal (maximal or minimal) on all plaquettes.

Stability depends on the eigenvalues of

$$\frac{\delta^2 S_{eff}}{\delta V_\mu(n_1) \delta V_\nu(n_2)} = \frac{\delta^2 S_{link}}{\delta V_\mu(n_1) \delta V_\nu(n_2)} + \frac{\delta^2 S_{plaq}}{\delta V_\mu(n_1) \delta V_\nu(n_2)}, \quad (10)$$

where

$$\frac{\delta^2 S_{link}}{\delta V \delta V} \sim \left(\frac{\beta}{4} \right)^{4(L-2)+12}, \quad \frac{\delta^2 S_{plaq}}{\delta V \delta V} \sim \left(\frac{\beta}{4} \right)^{L^2}. \quad (11)$$

Crucial fact is now the following: for $\beta/4 \ll 1$ and

$$4(L-2) + 12 < L^2 \quad \implies \quad L \geq 5 \quad (12)$$

the contribution of $\delta^2 S_{plaq} / \delta V \delta V$ to the stability matrix (and therefore to its eigenvalues) is negligible compared to $\delta^2 S_{link} / \delta V \delta V$.

This has (at least) three important implications:

- *Vortex stability*: The thin vortex is a stable saddlepoint of the full effective action S_{eff} at $L \geq 5$.
- *Vortex thickness*: A “thin” vortex on the V -lattice means thickness $< L$ on the U -lattice. This means that stable center vortices are $\approx 4-5$ lattice spacings thick. This is the distance where the adjoint string breaks at strong couplings!
- *Percolation*: From S_{eff} , we see that center vortices in $D = 3$ cost action $8(\beta/4)^{L^2}$ /unit length, while the entropy is $O(1)$ /unit length. Entropy \gg action implies that vortices percolate through the lattice, and confine nonzero N -ality charges.

However, the above result has been obtained in the strong-coupling approximation. To study the role of vortices at weak couplings, we resort to numerical simulations.

3 Center Projection in Maximal Center Gauge

The model of ‘t Hooft relies on no particular gauge. However, without gauge fixing, one can derive certain relations that, at the first sight, seem related to vortices, but apparently have no physical relevance¹⁸. Gauge fixing turns out to be also useful in pin-pointing the relevant (for vortices and confinement) degrees of freedom. One can then study in detail e.g. their topological properties¹³, behaviour across deconfinement phase transition¹², etc.

We proposed to identify vortices in thermalized lattice configurations in the following steps^{7,8}:

- First we fix to *maximal center gauge*^b by maximizing the expression:

$$\mathcal{R} \equiv \sum_{x,\mu} \left| \text{Tr}[U_\mu(x)] \right|^2. \quad (13)$$

This in fact is *adjoint Landau gauge*; the above condition is equivalent to maximizing

$$\mathcal{R}' \equiv \sum_{x,\mu} \text{Tr}[U_\mu^A(x)]. \quad (14)$$

- Then we make *center projection* by replacing:

$$U_\mu(x) \rightarrow Z_\mu(x) \equiv \text{sign Tr}[U_\mu(x)]. \quad (15)$$

- Finally we identify excitations (*P-vortices*) of the resulting Z_2 lattice configurations.

^bFor the continuum counterpart of MCG see the talk of Reinhardt^{1,19}.

A series of findings shows that vortices identified via center projection in MCG, are crucial for colour confinement and related phenomena:

- *Center dominance:* It was observed both in $SU(2)$ ^{7,8} and (less convincingly) in $SU(3)$ lattice gauge theory ²⁰ that the string tension obtained from center-projected configurations in MCG (or Laplacian center gauge ²¹) agrees remarkably well with the asymptotic string tension of the full theory.
- *P-vortices locate center vortices:* ⁸ One can define “vortex-limited” Wilson loops $W_n(C)$, i.e. Wilson loops evaluated on the original, unprojected lattice, on a subensemble of configurations in which n P-vortices pierce the minimal area of the loop. One expects $W_n(C)/W_0(C)$ to approach $(-1)^n$ asymptotically; this expectation is confirmed for large enough loops in numerical simulations.
- *P-vortices locate physical objects:* ^{10,8} Vortex density scales according to the asymptotic freedom, like a dimensionful quantity $\sim \Lambda^2$.
- *Center vortices are correlated with confinement:* ^{7,9} Removal of center vortices destroys confinement.
- *Center vortices are correlated with chiral symmetry breaking:* ⁹ Removal of center vortices restores chiral symmetry.
- *Center vortices cause non-trivial topology:* ⁹ Lattice configurations without vortices belong to the *trivial topological sector* (have zero topological charge).
- *Deconfinement can be understood as a center vortex percolation transition:* ^{12,13} In the confined phase, most P-plaquettes belong to a single huge vortex cluster percolating through the whole lattice volume. In the deconfined phase, vortex clusters cease percolating in *space* slices, while they continue percolating in *time* slices. Most vortices are short loops, winding in the time direction. In this way, perimeter-law behaviour of time-like Wilson loops can peacefully coexist with area-law behaviour of space-like Wilson loops.

4 Center Dominance, Gauge Copies, and Lattice Size

Though the above list of indications in favour of our vortex identification procedure looks quite impressive, a word of caution is necessary. MCG, similar to Landau or maximal abelian gauge, suffers from the *Gribov copy problem*. The iterative procedure, used for gauge fixing, converges to a local maximum which is slightly different for every gauge copy of a given lattice configuration. At the first sight, Gribov copies in MCG do not seem to be a severe problem in our procedure; it appears that P-vortex locations vary comparatively little,

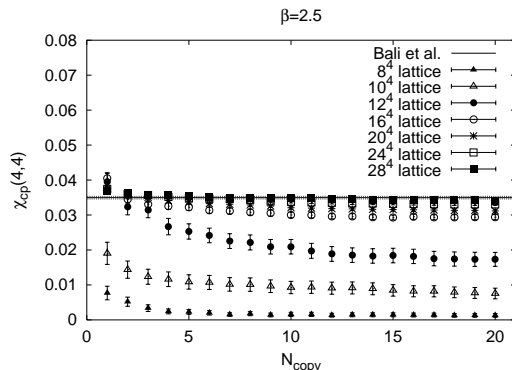


Figure 3: The dependence of the Creutz ratio $\chi(4, 4)$ on the number of gauge copies N_{copy} for $\beta = 2.5$, for various lattice sizes.

from copy to copy⁸.

However, a serious drawback might be a strong dependence of physical results (Creutz ratios, vortex density, etc.) on the number of gauge copies used in MCG maximization procedure. For each lattice configuration one can generate a set of its random copies, fix all these copies to MCG, and evaluate quantities of interest on the “best” copy, that is the one with the highest value of \mathcal{R} in Eq. (13). Such an investigation has recently been reported by Bornyakov et al.¹⁶ They e.g. show that the $N_{copy} \rightarrow \infty$ extrapolation of the center-projected string tension underestimates the full string tension by as much as 30% (at $\beta = 2.5$)! This seems to indicate that the observed center dominance in MCG (obtained from a small number of copies) might have been just a numerical coincidence.

In our opinion, the reported strong disagreement is due to finite size effects. The lattices used by Bornyakov et al.¹⁶ seem rather small compared to those used in our earlier studies. In fact, what is “rather small” or “large enough” depends on what are the typical sizes of vortex cores. Independent sources of estimating vortex thickness^{8,22,23} point towards a value of a little over one fermi, i.e. center vortices are ≈ 12 – 14 lattice spacings thick at $\beta = 2.5$ and the lattice size, used by Bornyakov et al. (16^4), seems too small. We have therefore repeated their analysis on lattices of various sizes; it turns out that on sufficiently large lattices the dependence on the number of gauge copies is much weaker, and center dominance is quite accurate. This fact is illustrated in Figs. 3 and 4, details can be found in our recent publication^{24,c}.

^cThis might not be the final answer to the problem. We have recently been informed that,

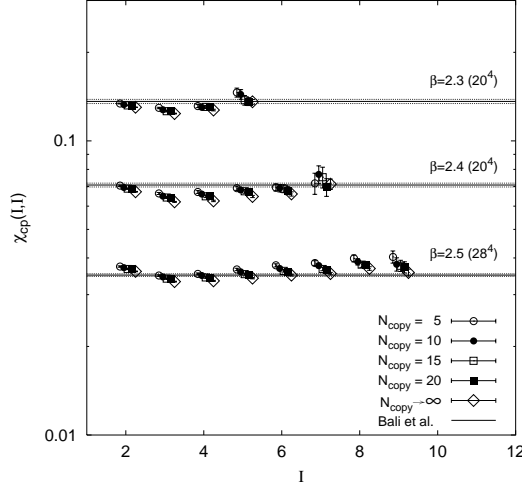


Figure 4: Center-projected Creutz ratios $\chi_{cp}(I, I)$ vs. I , from our largest lattices.

5 Casimir Scaling

We have repeatedly emphasized that trying to understand the confinement mechanism one has to be able to explain the problem of Casimir scaling. If we look at string tensions for higher representation colour charges, we find different behaviour in two distinguishable regimes:

- *Casimir-scaling regime:* At intermediate distances, from the onset of confinement to the onset of screening, the string tension is roughly proportional the quadratic Casimir. For $SU(2)$, e.g.,

$$\frac{\sigma_j}{\sigma_{1/2}} \approx \frac{4}{3}j(j+1). \quad (16)$$

- *Asymptotic regime:* The $SU(N)$ quark colour charge is eventually screened by gluons to the lowest representation with the same transformation

using a combination of simulated annealing and the usual (over)relaxation, one can still find copies with higher maxima of \mathcal{R} [V. Bornyakov, M. Polikarpov, private communication]. Also, using Landau gauge preconditioning, one reaches a higher maximum, but loses center dominance²⁵. Though we were able to find a cause of this problem, the loss of what we called “vortex-finding property”²⁶ during LG fixing, it is clear that a better gauge-fixing procedure and/or a modification of the gauge-fixing condition, is needed. Some alternatives have been proposed^{27,28,4}.

properties under Z_N . For $SU(2)$:

$$\sigma_j = \begin{cases} \sigma_{1/2} & \dots & j = \text{half-integer}, \\ 0 & \dots & j = \text{integer}. \end{cases} \quad (17)$$

As $N \rightarrow \infty$, colour screening is suppressed, Casimir scaling is exact. But even at $N = 2, 3$ there is a Casimir-scaling region of finite extent^{29,30}.

Center vortices are good at explaining the asymptotic region. Loops which transform trivially under the group center (such as the adjoint) are unaffected by fluctuations in the center elements, and therefore shouldn't get an area law. This is strictly true, however, only for *thin* center vortices (P-vortices).

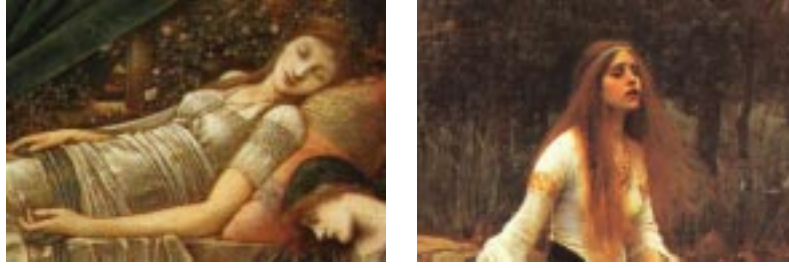
Is there a way of reconciling approximate Casimir scaling at intermediate distances with center vortices? The answer is apparently yes, and the crucial ingredient in the explanation of Casimir scaling through vortices is the vortex thickness. If the slice of a thick vortex, in the plane of a Wilson loop, is entirely inside (or entirely outside) the loop, the effect is the same as in the thin vortex model. However, if the vortex overlaps the loop perimeter, one can envisage its effect on the loop by insertion of a group element G , interpolating smoothly between -1 and 1 (as would be the case in an abelian theory). Were the vortices relatively thin on average, this effect would be unimportant; it would lead to perimeter-like correction to Wilson loops up to sizes comparable to the size of the vortex core. But, as already discussed in the previous section, there are good reasons to believe that vortices are quite thick. This enables one to explain the gross features of the behaviour of string tensions both in the Casimir-scaling and asymptotic regions. A simple model has been worked out in our earlier work³¹ (see also³²).

Bali²⁹ and Deldar³⁰ have recently published results of high statistics computations of higher representation potentials in the $SU(3)$ lattice gauge theory. Casimir scaling turns out to be quite precise, much more precise than in our over-simplified model. We believe this is not a real problem for center vortices, and our model and Ansatz for vortex-core profiles can be refined to explain the observed behaviour.

Finally I want to mention that to be compatible with $N \rightarrow \infty$ where Casimir scaling is exact, the vortex thickness should grow with N . Some evidence for such a growth has recently been presented by Montero³³.

6 Epilogue

The first in this series of Conferences took place in 1994 at the Lake of Como. There is no mention of the vortex model of confinement in its Proceedings. The idea, after its happy infant years, was sleeping deeply at that time like the



(a) E. Burne-Jones (1833–1898): *The Sleeping Princess (Briar Rose)*, detail. (b) J. W. Waterhouse (1849–1917): *The Lady of Shalott*, detail.

Figure 5: “Confinement I” (1994) versus “Confinement IV” (2000).

Briar Rose (Fig. 5a). I tried to convince you in my talk that now, six years later, the Sleeping Beauty is up after a long sleep, and looks very attractive (Fig. 5b). Some problems, however, persist. I sincerely hope that, in the years to come, she will not fall asleep again, or turn into the Maleficent Queen.

Acknowledgments

I am grateful to Wolfgang Lucha, Nora Brambilla, Khin Maung Maung and the other organizers for inviting me to present the talk at this stimulating meeting.

References

1. H. Reinhardt, these Proceedings.
2. A. Alexandru and R. W. Haymaker, hep-lat/0009010.
3. M. Engelhardt, H. Reinhardt and M. Faber, hep-lat/0008021.
4. K. Langfeld, hep-lat/0008019.
5. G. ’t Hooft, Nucl. Phys. **B138**, 1 (1978).
6. K. G. Wilson, Phys. Rev. **D10**, 2445 (1974).
7. L. Del Debbio, M. Faber, J. Greensite, and Š. Olejník, Phys. Rev. **D55**, 2298 (1997) [hep-lat/9610005].
8. L. Del Debbio, M. Faber, J. Giedt, J. Greensite, and Š. Olejník, Phys. Rev. **D58**, 094501 (1998) [hep-lat/9801027].
9. Ph. de Forcrand and M. D’Elia, Phys. Rev. Lett. **82**, 4582 (1999) [hep-lat/9901020].
10. K. Langfeld, H. Reinhardt and O. Tennert, Phys. Lett. **B419**, 317 (1998) [hep-lat/9710068].

11. K. Langfeld, O. Tennert, M. Engelhardt, and H. Reinhardt, Phys. Lett. **B452**, 301 (1999) [hep-lat/9805002].
12. M. Engelhardt, K. Langfeld, H. Reinhardt, and O. Tennert, Phys. Rev. **D61**, 054504 (2000) [hep-lat/9904004].
13. R. Bertle, M. Faber, J. Greensite, and Š. Olejník, J. High Energy Phys. **9903**, 019 (1999) [hep-lat/9903023].
14. B. L. Bakker, A. I. Veselov and M. A. Zubkov, Phys. Lett. **B471**, 214 (1999) [hep-lat/9902010];
M. N. Chernodub, M. I. Polikarpov, A. I. Veselov, and M. A. Zubkov, Nucl. Phys. Proc. Suppl. **73**, 575 (1999) [hep-lat/9809158].
15. M. Faber, J. Greensite, and Š. Olejník, J. High Energy Phys. **0006**, 041 (2000) [hep-lat/0005017].
16. V. G. Bornyakov, D. A. Komarov, M. I. Polikarpov, and A. I. Veselov, JETP Lett. **71**, 231 (2000) [hep-lat/0002017].
17. G. G. Batrouni, Nucl. Phys. **B208**, 12 (1982).
18. M. Faber, J. Greensite and Š. Olejník, J. High Energy Phys. **9901**, 008 (1999) [hep-lat/9810008].
19. M. Engelhardt and H. Reinhardt, Nucl. Phys. **B567**, 249 (2000) [hep-th/9907139].
20. M. Faber, J. Greensite and Š. Olejník, Phys. Lett. **B474**, 177 (2000) [hep-lat/9911006].
21. Ph. de Forcrand and M. Pepe, hep-lat/0008016.
22. Ph. de Forcrand and O. Philipsen, Phys. Lett. **B475**, 280 (2000) [hep-lat/9912050].
23. T. G. Kovács and E. T. Tomboulis, Phys. Rev. Lett. **85**, 704 (2000) [hep-lat/0002004].
24. R. Bertle, M. Faber, J. Greensite, and Š. Olejník, hep-lat/0007043.
25. T. G. Kovács and E. T. Tomboulis, Phys. Lett. **B463**, 104 (1999) [hep-lat/9905029].
26. M. Faber, J. Greensite, Š. Olejník, and D. Yamada, J. High Energy Phys. **9912**, 012 (1999) [hep-lat/9910033].
27. C. Alexandrou, M. D'Elia and Ph. de Forcrand, Nucl. Phys. Proc. Suppl. **83**, 437 (2000) [hep-lat/9907028].
28. B. L. Bakker, A. I. Veselov and M. A. Zubkov, hep-lat/0007022.
29. G. S. Bali, these Proceedings; hep-lat/0006022.
30. S. Deldar, Phys. Rev. **D62**, 034509 (2000) [hep-lat/9911008].
31. M. Faber, J. Greensite and Š. Olejník, Phys. Rev. **D57**, 2603 (1998) [hep-lat/9710039].
32. S. Deldar, hep-ph/9912428.
33. Á. Montero, Phys. Lett. **B483**, 309 (2000) [hep-lat/0004002].